

# Manifestly N=3 Supersymmetric Euler-Heisenberg Action in Light-Cone Superspace

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## Abstract

We find a manifestly N=3 supersymmetric generalization of the four-dimensional Euler-Heisenberg (four-derivative, or  $F^4$ ) part of the Born-Infeld action in light-cone gauge, by using N=3 light-cone superspace.

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# 1 Introduction

The *Born-Infeld* (BI) action in flat spacetime,<sup>2</sup>

$$S_{\text{BI}} = \frac{1}{b^2} \int d^4x \left\{ 1 - \sqrt{-\det(\eta_{\mu\nu} + bF_{\mu\nu})} \right\} , \quad (1.1)$$

is the particular non-linear generalization of Maxwell theory,  $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ . The action (1.1) was initially introduced to regularize both the electric field and the self-energy of a point-like charge in electrodynamics [1]. Much later, the BI action was recognized as the leading contribution to the effective action of open strings in an abelian background with constant field strength  $F$  [2], and as the essential part of the D3-brane action as well [3], with  $b = 2\pi\alpha'$ . The action (1.1) has many remarkable properties, e.g., causal propagation and electric-magnetic duality [4, 5].

The BI Lagrangian can be rewritten to the form

$$L = -\frac{1}{2}p^{\mu\nu}F_{\mu\nu} + H(P, Q) , \quad (1.2)$$

where the auxiliary antisymmetric tensor  $p_{\mu\nu}$  and the BI structure function

$$H(P, Q) = \frac{1}{b^2} \left( 1 - \sqrt{1 - 2b^2P + b^4Q^2} \right) , \quad (1.3)$$

as well as the definitions

$$P = \frac{1}{4}p_{\mu\nu}p^{\mu\nu} , \quad Q = \frac{i}{4}p_{\mu\nu}\tilde{p}^{\mu\nu} , \quad \tilde{p}^{\mu\nu} = \frac{1}{2}\varepsilon^{\mu\nu\rho\sigma}p_{\rho\sigma} , \quad (1.4)$$

have been introduced. Eliminating  $p_{\mu\nu}$  from eq. (1.2) results in the equivalent Lagrangian

$$L = \frac{1}{b^2} \left[ 1 - \sqrt{1 + \frac{b^2}{2}F^2 - \frac{b^4}{16}(F\tilde{F})^2} \right] , \quad (1.5)$$

where we have defined  $F^2 = F^{\mu\nu}F_{\mu\nu}$ ,  $\tilde{F}^{\mu\nu} = \frac{1}{2}\varepsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}$  and  $F\tilde{F} = F^{\mu\nu}\tilde{F}_{\mu\nu}$ .

Supersymmetric generalizations of the BI action are of particular interest in connection to superstring theory (see ref. [6] for a recent review). The super-BI actions describing D-branes can be naturally interpreted as the Goldstone-type actions associated with partial supersymmetry breaking, while they can still be duality invariant too. The manifestly N=1 supersymmetric generalization of the four-dimensional BI action in N=1 superspace was discovered long time ago [7] (see also ref. [8]), while its manifestly N=2 supersymmetric generalization in N=2 superspace was found only recently [9] (see ref. [10] too). To our knowledge, the higher (N=3 or N=4) manifestly

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<sup>2</sup>We use  $\eta_{\mu\nu} = \text{diag}(+, -, -, -)$  and  $\hbar = c = 1$ .

supersymmetric generalizations of the four-dimensional bosonic BI action (1.1) are not known in any form.

Supersymmetry apparently prefers the parametrization of the BI action in terms of the Maxwell term  $L_2 = -\frac{1}{4}F^2$  and the Maxwell stress-energy tensor squared [9],

$$L_4 = \frac{1}{32} \left\{ (F^2)^2 + (F\tilde{F})^2 \right\} = \frac{1}{8} (F^+)^2 (F^-)^2, \quad F_{\mu\nu}^\pm = \frac{1}{2} (F_{\mu\nu} \pm i\tilde{F}_{\mu\nu}) . \quad (1.6)$$

This term is known as the *Euler-Heisenberg* (EH) Lagrangian [11]. The EH action also appears as the bosonic part of the one-loop effective action in N=1 supersymmetric scalar electrodynamics with the parameter  $b^{-1} = 2\sqrt{6}\pi m^2/e^2$ . One easily finds that

$$L_{\text{BI}} = \frac{1}{b^2} \left\{ 1 - \sqrt{(1 - b^2 L_2)^2 - 2b^4 L_4} \right\} = L_2 + b^2 L_4 + O(F^6) . \quad (1.7)$$

A manifestly N=4 supersymmetric generalization of the BI action is known to be the formidable problem, though it is highly desirable, e.g., for an investigation of quantum properties of D3-branes and their comparison to supergravity [12, 13]. Even the  $N > 2$  supersymmetrization of the EH-term  $L_4$ , representing the four-derivative terms ( $F^4$ ), is non-trivial. The additional terms with four derivatives in the N=4 BI action were determined in ref. [14] in N=1 superspace, by imposing the  $SU(3)$  internal symmetry on three N=1 chiral multiplets extending an N=1 (abelian) vector multiplet to an N=4 vector multiplet, with manifest (linearly realised) N=1 off-shell supersymmetry. The manifestly N=2 supersymmetric form of the N=4 EH action was derived in ref. [15] in N=2 projective superspace, while its equations of motion can also be written in terms of on-shell N=4 superfields in harmonic superspace [16]. It is the purpose of this Letter to write down an off-shell, manifestly N=3 supersymmetric formulation of the N=4 EH action in N=3 light-cone superspace.

Our paper is organized as follows. In sect. 2 we introduce a light-cone gauge and rewrite the EH Lagrangian in terms of physical (transverse) degrees of freedom up to the relevant order. In sect. 3 we introduce N=3 light-cone superspace and deduce an N=3 supersymmetric generalization of the EH action in terms of a single N=3 light-cone superfield. The obstructions encountered in our efforts to find a similar, manifestly N=4 supersymmetric EH action in N=4 light-cone superspace are discussed in Conclusion (sect. 4).

## 2 EH action in light-cone gauge

The light-cone formulation of a gauge theory (in light-cone gauge) keeps only physical (transverse) degrees of freedom in the field theory by giving up its manifest

Lorentz invariance. The light-cone formulation is, therefore, very suitable for an off-shell formulation of N-extended supersymmetric gauge field theories with manifest supersymmetry in N-extended light-cone superspace [17, 18, 19].

We define light-cone coordinates in Minkowski spacetime as

$$\begin{aligned} x^+ &= \frac{1}{\sqrt{2}} (x^0 + x^3) , & x^- &= \frac{1}{\sqrt{2}} (x^0 - x^3) , \\ x &= \frac{1}{\sqrt{2}} (x^1 + ix^2) , & \bar{x} &= \frac{1}{\sqrt{2}} (x^1 - ix^2) , \end{aligned} \quad (2.1)$$

and similarly for the gauge vector field,  $A_\mu \rightarrow (A^+, A^-, A, \bar{A})$ . The real coordinate  $x^+$  is going to be considered as ‘light-cone time’. The linear transformation (2.1) of spacetime coordinates is obviously non-singular (with the Jacobian equal to  $i$ ), while it does not preserve the Minkowski metric (i.e. it is not a Lorentz-transformation).

The light-cone gauge reads

$$A^+ = 0 . \quad (2.2)$$

In this (physical) gauge the  $A^-$  component of the gauge field  $A_\mu$  is supposed to be eliminated via its (non-dynamical) equation of motion, whereas the transverse components  $(A, \bar{A})$  are supposed to represent the physical propagating fields.

It is easy to solve the equation of motion for  $A^-$  in the Maxwell theory, where it takes the form of a linear equation in the light-cone gauge (*cf.* refs. [17, 18, 19]). It becomes, however, a highly non-trivial problem in the BI or EH theory, where it takes the form of a non-linear partial differential equation. The equations of motion amount to the conservation law for the  $p$ -tensor,

$$\partial^\mu p_{\mu\nu} = 0 , \quad (2.3)$$

while the  $p_{\mu\nu}$  in the BI theory is given by

$$p_{\mu\nu} = \frac{b^2 F_{\mu\nu} - \frac{b^4}{4} (F\tilde{F})\tilde{F}_{\mu\nu}}{\sqrt{1 + \frac{b^2}{2} F^2 - \frac{b^4}{16} (F\tilde{F})^2}} . \quad (2.4)$$

By the use of the Bianchi identity,  $\partial^\mu \tilde{F}_{\mu\nu} = 0$ , we find the following equation for  $A^-$ :

$$\begin{aligned} \partial^\mu F_{\mu-} &= b^2 \left\{ -\frac{1}{2} \partial^\mu F_{\mu-} F^2 + \frac{1}{4} \tilde{F}_{\mu-} \partial^\mu (F\tilde{F}) + \frac{1}{4} F_{\mu-} \partial^\mu F^2 \right\} \\ &+ b^4 \left\{ -\frac{1}{16} (F\tilde{F}) \tilde{F}_{\mu-} \partial^\mu F^2 + \frac{1}{16} \partial^\mu F_{\mu-} (F\tilde{F})^2 \right. \\ &\left. + \frac{1}{16} \tilde{F}_{\mu-} \partial^\mu (F\tilde{F}) F^2 - \frac{1}{32} F_{\mu-} \partial^\mu (F\tilde{F})^2 \right\} . \end{aligned} \quad (2.5)$$

We use a perturbative *Ansatz*, in powers of the small parameter  $b^2$ , for a solution to eq. (2.5),

$$A^-(x) = \sum_{n=0}^{\infty} b^{2n} A_{(2n)}^-(x) . \quad (2.6)$$

As regards the leading and sub-leading terms, we find

$$\begin{aligned} A_{(0)}^- &= \frac{1}{\partial^+} (\bar{\partial} A + \partial \bar{A}) , \\ A_{(2)}^- &= \frac{1}{(\partial^+)^2} \left[ -\frac{1}{2} \partial^\mu F_{\mu-} F^2 + \frac{1}{4} \tilde{F}_{\mu-} \partial^\mu (F \tilde{F}) + \frac{1}{4} F_{\mu-} \partial^\mu F^2 \right] \Big|_{A^- = A_{(0)}^-} , \end{aligned} \quad (2.7)$$

where we have used the notation  $\partial^+ = \partial/\partial x^-$ . The multiple factors  $(\partial^+)^{-1}$  in our actions are harmless after rewriting them to momentum space. The first line of eq. (2.7) coincides with the exact solution in the Maxwell theory.

According to eq. (1.7), the EH term  $L_4$  is the leading  $b^2$ -correction to the Maxwell term  $L_2$  in the BI theory. The light-cone formulation of the BI Lagrangian in the same approximation is thus given by the terms written down on the right-hand-side of eq. (1.7) after a substitution of eq. (2.2) and the first line of eq. (2.8). After some algebra and partial integration we find

$$\begin{aligned} L[A, \bar{A}] &= -\frac{1}{4} F^2 + \frac{b^2}{8} (F^+)^2 (F^-)^2 \\ &= -A \square \bar{A} + 2b^2 \left| (\partial \bar{A})^2 + \partial^+ \bar{A} \frac{\square}{2\partial^+} A - \partial^+ \bar{A} \frac{\partial^2}{\partial^+} \bar{A} \right|^2 + O(b^4) , \end{aligned} \quad (2.8)$$

where we have used the notation  $\partial = \partial/\partial x$  and  $\bar{\partial} = \partial/\partial \bar{x}$ . Eq. (2.8) can be thought of as the light-cone EH Lagrangian. Its N=3 supersymmetrization is discussed in the next sect. 3.

### 3 N=3 light-cone superspace action

The light-cone N=3 supersymmetry algebra reads

$$\{Q^m, \bar{Q}_n\} = -\sqrt{2} \delta_n^m P^+ , \quad m, n = 1, 2, 3 , \quad (3.1)$$

where the supersymmetry charges  $Q^r$  transform in the fundamental representation of  $SU(3)$ . A natural representation of the algebra (3.1) in N=3 light-cone superspace  $Z = (x^\mu, \theta^m, \bar{\theta}_n)$  is given by

$$\begin{aligned} Q^m &= -\frac{\partial}{\partial \theta_m} + \frac{i}{\sqrt{2}} \theta^m \partial^+ , \\ \bar{Q}_n &= \frac{\partial}{\partial \theta^n} - \frac{i}{\sqrt{2}} \bar{\theta}_n \partial^+ . \end{aligned} \quad (3.2)$$

The covariant derivatives in N=3 light-cone superspace are

$$\begin{aligned} D^m &= -\frac{\partial}{\partial \bar{\theta}_m} - \frac{i}{\sqrt{2}} \theta^m \partial^+ , \\ \bar{D}_n &= \frac{\partial}{\partial \theta^n} + \frac{i}{\sqrt{2}} \bar{\theta}_n \partial^+ . \end{aligned} \quad (3.3)$$

They anticommute with the supersymmetry charges (3.2) and obey the same algebra (3.1). The irreducible off-shell representations of N=3 light-cone supersymmetry are easily obtained by imposing the covariant chirality condition on N=3 light-cone superfields  $\phi(Z)$ ,

$$D^m \phi(Z) = 0 . \quad (3.4)$$

A solution to eq. (3.4) in components is just given by an arbitrary complex function  $\phi(x^+, x^- + \frac{i}{\sqrt{2}} \theta^m \bar{\theta}_m, x, \bar{x}; \theta^n) \equiv \phi(y; \theta)$ . Its expansion in the chiral superspace reads

$$\phi(y; \theta) = \frac{1}{\partial^+} A(y) + \frac{i}{\partial^+} \theta^m \bar{\chi}_m(y) + \frac{i}{2} \theta^m \theta^n \varepsilon_{mnp} C^p(y) + \frac{1}{3!} \varepsilon_{mnp} \theta^m \theta^n \theta^p \psi(y) . \quad (3.5)$$

The light-cone N=3 supersymmetry transformation laws for the components are

$$\begin{aligned} \delta A &= i \varepsilon^n \bar{\chi}_n , \\ \delta \bar{\chi}_m &= \sqrt{2} \bar{\varepsilon}_m \partial^+ A + \varepsilon_{mnp} \varepsilon^n \partial^+ C^p , \\ \delta C^p &= -i \sqrt{2} \varepsilon^{pqr} \bar{\varepsilon}_q \bar{\chi}_r - i \varepsilon^p \psi , \\ \delta \psi &= -\sqrt{2} \bar{\varepsilon}_n \partial^+ C^n , \end{aligned} \quad (3.6)$$

where  $(\varepsilon^n, \bar{\varepsilon}_m)$  are the infinitesimal anticommuting parameters.

All our field components have canonical dimensions. The complex field  $A$  can be identified with the physical (translational) vector field components, the spinors  $\bar{\chi}_m$  in the fundamental representation **3** of  $SU(3)$  with a triplet of photinos, the singlet spinor  $\psi$  with extra photino, and the complex triplet  $C^m$  with Higgs fields in **3** of  $SU(3)$ . The physical content thus coincides with that of the N=4 supersymmetric abelian vector multiplet having a single photon field, photinos in the fundamental representation **4** of  $SU(4)$  and Higgs fields in real **6** of  $SU(4)$ , after their decomposition with respect to the  $SU(3)$  subgroup of the internal symmetry  $SU(4)$ . This is the manifestation of the well-known fact that N=3 and N=4 supersymmetric vector multiplets are physically equivalent.

It is now straightforward (though very tedious) to find the N=3 supersymmetric generalization of the bosonic EH light-cone action (2.8) in N=3 light-cone superspace,

$$S = \int d^4 x d^3 \theta d^3 \bar{\theta} \mathcal{L}(\phi, \bar{\phi}) = - \int d^4 x (D)^3 (\bar{D})^3 \mathcal{L}(\phi, \bar{\phi}) , \quad (3.7)$$

where  $(D)^3 = \varepsilon_{mnp} D^m D^n D^p$  and similarly for  $(\bar{D})^3$ . After some trials and errors, we find

$$\begin{aligned}
36(-i\sqrt{2})^3 \mathcal{L}(\phi, \bar{\phi}) = & -\phi \frac{\square}{\partial^+} \bar{\phi} + 2b^2 \left\{ \frac{1}{\partial^{+3}} (\bar{\partial} \partial^+ \phi \bar{\partial} \partial^+ \phi) (\partial \partial^+ \bar{\phi})^2 \right. \\
& + \frac{1}{\partial^{+3}} (\partial^{+2} \phi \bar{\partial}^2 \phi) \partial^{+2} \bar{\phi} \partial^2 \bar{\phi} + \frac{1}{2\partial^+} (\phi) (\partial \partial^+ \bar{\phi})^2 \square \bar{\phi} \\
& + \frac{1}{4\partial^{+3}} (\partial^{+2} \phi \square \phi) \partial^{+2} \bar{\phi} \square \bar{\phi} - \frac{1}{\partial^{+3}} (\partial^{+2} \phi \bar{\partial}^2 \phi) (\partial \partial^+ \bar{\phi})^2 \\
& - \frac{1}{2\partial^{+3}} (\partial^{+2} \phi \square \phi \bar{\partial}^2 \phi) \partial^{+2} \bar{\phi} - \frac{1}{\partial^{+3}} (\bar{\partial} \partial^+ \phi \bar{\partial} \partial^+ \phi) \partial^{+2} \bar{\phi} \partial^2 \bar{\phi} \\
& \left. + \frac{1}{2\partial^{+3}} (\bar{\partial} \partial^+ \phi \bar{\partial} \partial^+ \phi \square \phi) \partial^{+2} \bar{\phi} - \frac{1}{2\partial^+} (\phi) \partial^{+2} \bar{\phi} \square \bar{\phi} \partial^2 \bar{\phi} \right\} . \tag{3.8}
\end{aligned}$$

The bosonic part of this action is given by

$$\begin{aligned}
\mathcal{L}_{\text{bos.}} = & -A \square \bar{A} + 2b^2 \left| (\partial \bar{A})^2 + \partial^+ \bar{A} \frac{\square}{2\partial^+} A - \partial^+ \bar{A} \frac{\partial^2}{\partial^+} \bar{A} \right|^2 \\
& - \frac{1}{2} C^p \square \bar{C}_p - 2b^2 \left\{ \frac{2}{\partial^{+2}} (\bar{\partial} \partial^+ C^m \bar{\partial} \partial^+ A) (\partial \partial^+ \bar{C}_m \partial \partial^+ \bar{A}) \right. \\
& + \frac{1}{2\partial^{+2}} (\partial^{+2} C^p \bar{\partial}^2 A + \bar{\partial}^2 C^p \partial^{+2} A) (\partial^{+2} \bar{C}_p \partial^2 \bar{A} + \partial^2 \bar{C}_p \partial^{+2} \bar{A}) \\
& + \frac{1}{8\partial^{+2}} (\partial^{+2} C^p \square A + \square C^p \partial^{+2} A) (\partial^{+2} \bar{C}_p \square \bar{A} + \square \bar{C}_p \partial^{+2} \bar{A}) \\
& + \left[ \frac{1}{4} C^p (2\partial \partial^+ \bar{C}_p \partial \partial^+ \bar{A} \square \bar{A} + \square \bar{C}_p \partial \partial^+ \bar{A} \partial \partial^+ \bar{A}) \right. \\
& - \frac{1}{\partial^{+2}} (\partial^{+2} C^p \bar{\partial}^2 A + \bar{\partial}^2 C^p \partial^{+2} A) \partial \partial^+ \bar{C}_p \partial \partial^+ \bar{A} \\
& \left. \left. - \frac{1}{4\partial^{+2}} (\partial^{+2} C^p \square A \bar{\partial}^2 A + \bar{\partial}^2 C^p \partial^{+2} A \square A + \square C^p \bar{\partial}^2 A \partial^{+2} A) \partial^{+2} \bar{C}_p + \text{h.c.} \right] \right\} . \tag{3.9}
\end{aligned}$$

One of the obvious features of both eqs. (3.8) and (3.9) is the apparent presence of higher derivatives, as may have been expected from the experience with the manifestly N=2 supersymmetric generalization of the BI action in the covariant N=2 superspace [9]. The expected correspondence to the component D3-brane effective action having non-linearly realized extended supersymmetry and no higher derivatives implies the existence of a field redefinition that would eliminate the higher-derivative terms in our action and make its N=3 supersymmetry to be non-linearly realised (i.e non-manifest) [6]. We also note the absence of quartic ( $C^4$ ) scalar terms and the on-shell ( $\square A = \square C = 0$ ) invariance of our action under constant shifts,  $C_p(x) \rightarrow C_p(x) + c_p$ , which are supposed to be related to the possible interpretation of the  $C_p$  fields as the Goldstone scalars associated with spontaneously broken translations in the full N=3 BI action.

## 4 Conclusion

Our main results are given by eqs. (2.8), (3.8) and (3.9). Our initial motivation was to construct an N=4 supersymmetric generalization of the EH action in the light-cone gauge. The N=4 light-cone supersymmetry algebra is given by eq. (3.1), where the indices now take four values. Equations (3.2), (3.3) and (3.4) are still valid in N=4 light-cone superspace, where they have to be supplemented by an extra (generalized reality) condition [17],

$$D^m D^n \bar{\phi} = \frac{1}{2} \varepsilon^{mnpq} \bar{D}^p \bar{D}_q \phi, \text{ or, equivalently, } \bar{\phi} = \frac{1}{48 \partial^+} \varepsilon^{mnpq} \bar{D}_m \bar{D}_n \bar{D}_p \bar{D}_q \phi. \quad (4.1)$$

The restricted chiral N=4 light-cone superfield  $\phi$  is equivalent to the chiral N=3 superfield in eq. (3.5). Our efforts to construct an N=4 generalization of eq. (2.8) along the similar lines (sect. 3) unexpectedly failed, while eq. (4.1) was the main obstruction. We conclude that even a manifestly N=4 supersymmetric generalization of the EH action in the light-cone gauge seems to be highly non-trivial, if any, not to mention an even more ambitious (manifest) N=4 supersymmetrization of the BI action.

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